

The Quantum-Classical Metal

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Abstract

In a normal Fermi liquid, Landau's theory precludes the loss of single fermion, quantum coherence in the low energy/temperature limit. For highly anisotropic, strongly correlated metals there is no proof that this remains the case: we propose that quantum coherence for transport in some directions may be lost intrinsically. This should stabilize a novel, qualitatively anisotropic non-Fermi liquid, separated by a novel zero temperature, quantum phase transition from the Fermi liquid state and categorized by the unobservability of certain interference effects. There is compelling experimental evidence for this transition as a function of magnetic field in the metallic phase of the organic conductor $(\text{TMTSF})_2\text{PF}_6$.

I. INTRODUCTION

For decades, the accepted theory of metallic states of electronic matter has been Landau's Fermi liquid theory (FLT). Recently, in the context of systems with strong electronic correlations, the universality of FLT has come into question. It is generally agreed that correlations are most dangerous to FLT when the electrons live in low (*i.e.* one or two) dimensions, and thus recent interest in breakdowns of FLT has focused on systems, such as the cuprate and organic superconductors (above their respective superconducting transition temperatures or in strong magnetic fields), which are “metallic”, bulk crystals, but enormously anisotropic. In these systems, the bare Coulomb interaction is not small, and as such it is not *a priori* obvious that it can be dealt with by perturbing about an anisotropic three-dimensional (3D) Fermi gas, as FLT requires. It is more natural to think of such systems as collections of strongly correlated one or two dimensional electron liquids with weak, interliquid single particle hopping. Since an essential characteristic of a Fermi liquid is the coherence of single electron motion in *all* spatial directions, the question arises as to whether such coherence is guaranteed to occur in the presence of anisotropy and strong interaction. The purpose of this paper is to introduce the reader to both theoretical and experimental considerations strongly challenging the universality and inevitability of such coherence. In fact, we propose here the existence of a new metallic, non-Fermi liquid state in which 3D quantum coherence is intrinsically destroyed by strong correlations *at zero temperature and in a pure crystal*. To orient the reader, we begin with a discussion of quantum coherence and how such coherence can be lost, leading to incoherent, or “classical”, behavior.

II. DECOHERENCE AND THE QUANTUM-CLASSICAL “BOUNDARY”

According to the Copenhagen interpretation of quantum mechanics there exists a remarkable, qualitative distinction between the “microscopic” and “macroscopic” realms. The microscopic world is governed by quantum mechanics, the macroscopic world by classical

mechanics; “observers” and “measuring apparatus” *etc.* lie in the latter, and “measurement” and the “collapse of the wavefunction” provide the procedure for describing the interactions of the micro- and macro-worlds. However, within the Copenhagen interpretation, no explicit means of determining the scale at which things become classical is given, and the means by which classical behavior emerges at the macroscale is not considered: it may or may not require new physics beyond the quantum rules that govern the micro-world. There is, however, no experimental evidence for a breakdown of quantum mechanics on any scale [1]; rather, classical behavior emerges in situations where either (1) the system is too complex for the quantum theory to make testable predictions, or (2) the chief hallmark of classical behavior - namely, the unobservability of interference effects between different histories - is in agreement with the predictions of quantum theory *when one considers the effects of the “environment” (the unobserved degrees of freedom in the experiment) on the magnitude of these interference effects.*

Recall that, in the Schrödinger picture of quantum mechanics, if a system is prepared in an energy eigenstate Ψ_E , then under time evolution the state picks up a temporal phase $-Et/\hbar$, *i.e.*, $\Psi_E \rightarrow e^{-iEt/\hbar}\Psi_E$. If a system is prepared in a superposition of such states, then under time evolution each eigenstate develops a phase independently of the others. If a subsequently applied perturbation allows the mixing of these states at a later time, then it is *in principle* possible to observe the spectacular oscillation effects of quantum mechanics. One of the more celebrated of these effects is strangeness oscillations in Kaon decay, but it is also perfectly correct to consider the ballistic motion of an electron wavepacket in an ordinary metal, as a true manifestation of coherent interference effects. In these cases, the interference exists because the energies of the states mixed by the relevant “perturbations” (the weak interaction and the intersite hopping, respectively) are very narrowly distributed with respect to the Hamiltonians in the absence of the perturbations. If one dealt instead with states which had energies specified only to within some range ΔE (and the eigenstates of the system are dense within this range) then the phase information would be lost in a time, $\Delta t \sim \hbar/\Delta E$, and Kaon oscillations or any other quantum interference effects would only be

observable on time scales shorter than or comparable to Δt . If the perturbation considered is too weak to have significant effects on this time scale, then no interference effects are observed and, further, discussing phase coherent superpositions of the mixed states is meaningless because the dephasing eliminates all the potentially observable coherence effects.

According to the decohering-histories interpretation of quantum mechanics [1], this loss of phase coherence is behind our generic inability to observe quantum coherence effects in the macro-world: macroscopic systems do obey quantum mechanics; however, in a macroscopic system it is generically a hopeless exercise to (1) specify all of the necessary degrees of freedom to restrict the system to a sufficiently narrow range of initial energies, and/or (2) find a perturbation which connects an initial state only to states which are themselves sufficiently degenerate in energy for observing quantum effects. In general, one can only restrict the wavefunction to a manifold in Hilbert space and experimentally realizable perturbations connect each state on this manifold to states which lie on some other manifold in Hilbert space. The energy of the states on these manifolds, *i.e.* the energy of the “initial” and/or “final” states, is defined only to within some spread ΔE , and rapid dephasing is the result. From this point of view, the Schrödinger’s cat *gedanken* experiment is spurious: it is meaningless to form superpositions of $|\text{dead}\rangle$ and $|\text{alive}\rangle$ states of the cat, since one cannot place the cat (even approximately) in an initial energy eigenstate nor find a perturbation which transforms the cat from an “alive” eigenstate to a group of sufficiently degenerate “dead” eigenstates to perform any kind of interference experiment. It may be unsettling that within the Copenhagen interpretation there is no deterministic way to know whether the cat will be alive or dead upon opening the box, but the more unsettling concept of the cat being in a superposition of $|\text{dead}\rangle$ and $|\text{alive}\rangle$ is meaningless in the sense of having no observable consequences [2].

III. DECOHERENCE IN “STRANGE” METALS

Interestingly, the modern theory of clean metals, Fermi liquid theory [3], is an exception to many of the generic statements made above: it predicts observable interference effects for certain *macroscopic* quantities. This is of particular interest because of the experimental evidence now accumulating for the existence of “strange” metals, *i.e.* materials whose properties are metallic, but which disobey various predictions of FLT. In this case, one might expect that some or all macroscopic quantum coherence which occurs for Fermi liquids might be removed for these strange metals. We will argue here that this is the case, that there exist strange metallic states in which the transport in one or more directions is “incoherent” in the sense that interference effects between histories which involve the transport of electrons in such direction(s) are unobservable *even in principle* (in the low energy limit for pure systems). This loss of coherence has strong physical consequences, which we will argue *have been experimentally observed*, demonstrating the existence of a new state of matter which might best be described as a “quantum-classical metal” because of the lack of quantum coherence in some direction(s).

Such a state is intrinsically and qualitatively anisotropic and could only be expected to arise in a material with strong anisotropy (and strong electron-electron interactions). In fact, our proposal arose out of a careful study of the problem of one dimensional chains of strongly interacting electrons, coupled by a very weak interchain single electron hopping, t_{\perp} . It is useful for orientation to briefly describe that study.

We begin with the observation that for short ranged, repulsive interactions, electrons confined to one dimension form not Fermi liquids but rather *Luttinger liquids* [4], a specific kind of “non-Fermi liquid”. The low energy properties of Luttinger liquids are well understood and there is therefore a firm basis for studying the action of t_{\perp} perturbatively [5,6]. It has long been known that t_{\perp} is the most relevant inter-liquid operator for a range of interaction strengths, implying that for such interactions, the correct low energy description of the system requires that t_{\perp} be retained regardless of how small it is. However, various

earlier studies of the problem also either suggested or assumed implicitly that the leading relevance of t_{\perp} was sufficient to establish the existence of *coherent*, interchain motion. The proposition is not self-evident and it was an attempt to critically address it for the coupled chains problem [5,6] that led to our proposal of the quantum-classical metal.

For coupled Luttinger liquids, there are a range of possible probes of interchain interference effects including, but not limited to, the frequency dependent transverse conductivity, the shape, if any, of the Fermi surface, the transverse bandwidth, the single particle Green's function, *etc.* In the quantum-classical metal, supposed to occur for sufficiently strong intrachain interactions, all of these probes of interliquid coherence should demonstrate the unobservability of interference effects between different histories involving interchain hops. Our original choice for a natural probe was provided by the following quantum oscillation effect: consider two identical chains of strongly interacting electrons, both in their isolated chain groundstates at time $t = 0$, but with chain 1 having δN more electrons than chain 2. At time $t = 0$, turn on the interchain hopping, t_{\perp} , and study the behavior of $\langle \delta N(t) \rangle$. This choice is motivated by the close connection between $\langle \delta N(t) \rangle$ and a quantity, $\langle \sigma^z(t) \rangle$ [5,6], extensively studied in the simplest model for the quantum to classical crossover, the Caldeira-Leggett model (hereafter the CL model) [7].

IV. ELECTRON HOPPING BETWEEN NON-FERMI LIQUIDS: CONNECTION TO A SIMPLE MODEL OF THE QUANTUM-CLASSICAL TRANSITION

In the CL model, σ represents the two state “macroscopic” degree of freedom (DOF) which is coupled to an environment of infinitely many microscopic degrees of freedom, represented by the simplest possible realization: a bath of harmonic oscillators. The oscillators “measure” the σ^z state of the DOF, while an additional perturbation (proportional to σ^x) mixes the two σ^z states coherently. The coupling to the oscillators should decohere the macroscopic degree of freedom (*i.e.* make superpositions of σ^z states meaningless) resulting in “classical behavior” for strong enough coupling to the oscillators. The quantity most

frequently studied is the expectation value, $\langle \sigma^z(t) \rangle$, where the system has been prepared by clamping σ^z to +1 for all $t < 0$ and allowing the environment to adapt to this configuration. The basic idea is that this resembles an experimentally realizable situation: the controllable, macroscopic degree of freedom is held in a particular state and the microscopic degrees of freedom, uncontrolled by the experimenter, relax to their equilibrium under these circumstances. The system is then released and one looks for quantum interference effects in the ensuing behavior of the observable, macroscopic degree of freedom.

One can also make a canonical transformation in the CL problem, changing basis to the eigenstates of the joint oscillator-DOF system (in the absence of the σ^x perturbation which mixes the two σ^z eigenstates). The CL Hamiltonian then takes the form

$$H_{\text{CL}} = \frac{1}{2} \Delta (\sigma^+ e^{-i\Omega} + \text{h.c.}) + \sum_i \left(\frac{1}{2} m_i \omega_i x_i^2 + \frac{1}{2m_i} p_i^2 \right)$$

where C_i is the coupling to the i th harmonic oscillator, m_i , ω_i , x_i and p_i are the mass, frequency, position and momentum of the i th oscillator, and $\Omega = \sum_i \frac{C_i}{m_i \omega_i} p_i$. In this language the “measurement” effects of the environment are encoded in the *non-degeneracy* of the states which are connected by the Δ term. In the absence of coupling to the bath (*i.e.* all C_i zero), the Δ term connects two *degenerate* σ_z states. In the presence of the bath, however, decoherence results when this degeneracy is sufficiently reduced by the operator $e^{\pm i\Omega}$ which creates and destroys oscillator bosons over a broad energy range whenever a transition between the σ^z states takes place.

In the new basis the usual CL preparation amounts to taking the system to be in one of the two groundstates of the system in the absence of the Δ term, and then suddenly switching this term on at time $t = 0$. This is parallel to the preparation discussed above for coupled Luttinger liquids; σ^z plays the role of δN , Δ the role of t_\perp and the oscillator bath the role of the well known charge and spin density oscillator modes of the coupled Luttinger liquids [4]. In both cases one follows the dynamics of the expectation value of a discrete, observable “macroscopic” variable (either σ^z or δN) which has been set up in a non-equilibrium state defined as a valid groundstate of the problem in the absence of

the perturbation (either the σ^x term or the inter-liquid single particle hopping term). As for the Δ term in the TLS, the action of the single particle hopping can be written as the product of an operator whose only action is to change δN , and an exponential in the creation and annihilation operators of the charge and spin density oscillator modes of the coupled Luttinger liquids, and the resulting Hamiltonian is strikingly similar to the CL Hamiltonian, H_{CL} [6]. Despite the fact that the models cannot be mapped into one another (*e.g.* the δN variable is many, not two, valued) the analogy does motivate the proposal that incoherent dynamics can occur for δN in a similar manner to the incoherence occurring for σ^z [6].

The signature of quantum coherence in the CL problem is taken to be the presence of oscillations in $\langle \sigma^z(t) \rangle$, which, when present, result from interference between histories in which σ^z varies differently. We correspondingly take the presence of oscillations in $\langle \delta N(t) \rangle$, as the signature of quantum coherence for the interchain hopping.

In FLT, $\langle \delta N \rangle$ exhibits oscillations with frequency Zt_{\perp} and damping which vanishes in the limit of vanishing t_{\perp} and $\delta N(t=0)/L$ [8]. FLT therefore exhibits a dramatic instance of macroscopic quantum coherence: not only are the oscillations observable, they are essentially undamped. The source of this can be traced back to the fact that, even though $\langle \delta N \rangle$ represents a macroscopic variable which is *a priori* expected to couple to a huge number of uncontrolled (and potentially dephasing) microscopic degrees of freedom, FLT dictates that $\langle \delta N \rangle$ is determined by the sum of many independent, coherent, quasiparticle channels. Since each channel decouples from its environment in the limit of vanishing $\delta N(t=0)/L$, coherence is unavoidable. However, for non-Fermi liquids, such as the Luttinger liquids of the coupled chains problem, there is no such special protection for the coherence of $\langle \delta N \rangle$.

In fact, while FLT is analogous to the CL model with no coupling between the oscillators and σ^z , coupled Luttinger liquids are analogous to the CL model with finite coupling to an ohmic bath of oscillators [5,6]. As a result, the most likely behaviors for $\langle \delta N \rangle$ fall into three categories: (1) for weak interactions, coherence and the characteristic oscillations, (2) for very strong interactions, localization with $\langle \delta N(t \rightarrow \infty) \rangle \neq 0$ and (3) for a range of intermediate interactions, incoherence, with no oscillations in $\langle \delta N \rangle$ but $\langle \delta N(t \rightarrow \infty) \rangle = 0$.

A remarkable result in the CL problem is that the final possibility is believed to occur over a broad range of coupling constants, so that it is *not* the case that the oscillations simply become more heavily damped as the localization behavior is approached. Rather, the oscillation frequency vanishes at some intermediate coupling before the localization sets in. The physical ingredients for this behavior are also present in the coupled Luttinger liquid problem [6] and we believe that case (3) should occur there as well. If the oscillation frequency of $\langle\delta N\rangle$ is identically zero (over some range of couplings where t_\perp is the leading instability of the uncoupled chains fixed point), then it is plausible that *all* interchain interference effects are unobservable in the low energy limit for this range of couplings. After all, there is nothing special about $\langle\delta N\rangle$. In the CL context, it is generally believed that the lack of coherence in $\langle\sigma^z\rangle$ signals a general loss of quantum coherence, and, in fact, the model was intended in part to explain the absence of interference effects for macroscopic objects. This absence is manifestly generic, rather than being limited to particular experiments. It is therefore quite likely that the disappearance of interference oscillations in $\langle\delta N\rangle$ represents a generic loss of coherence, rather than one limited to the probe considered. We correspondingly expect that in this regime the transverse electrical conductivity should lack a Drude peak, the single particle Greens function should not exhibit a pole on the real axis which disperses with k_\perp [6], no pair of split Fermi surfaces should form, *etc.*

For two chains in the incoherent phase, we expect there to be no Fermi surface splitting, which translates for infinitely many coupled chains into the absence of warping of any higher dimensional Fermi surface. This implies that, if it exists, this regime constitutes a new state of matter as the Fermi surface shape gives a clear, zero temperature, infinite time distinction between the incoherent phase and a normal metal. We believe that the incoherent state is separated by a zero temperature quantum phase transition from a state where the interchain hopping is coherent and a three dimensional, Fermi liquid metal occurs.

At this point it is important to emphasize that our proposed phase is *not* one in which the electrons are confined to the chains (this would be the analog of (2) above for the CL model). The latter is a phase in which, in the language of the renormalization group, t_\perp is

an *irrelevant* operator. Our proposed phase is truly novel in that it is only the coherence which is confined, not the electrons - diffusive interchain motion still takes place, and t_{\perp} is a *relevant* operator.

Our proposal is not limited to the case of coupled Luttinger liquid chains. While theoretically this was the best controlled case to study and the one most closely related to the CL model, there are other possibilities. In particular, a set of strongly interacting two dimensional systems, whose isolated groundstates were non-Fermi liquid metals would be natural candidates for incoherent interplane hopping. It is this possibility which we will now argue is experimentally realized in the organic conductor $(\text{TMTSF})_2\text{PF}_6$.

V. THE QUANTUM-CLASSICAL METALLIC STATE IN $(\text{TMTSF})_2\text{PF}_6$

$(\text{TMTSF})_2\text{PF}_6$ is a Bechgaard salt composed of linear stacks of tetramethyltetraselenafulvalene cations; the stacks are arranged into planes separated by PF_6 anions which provide overall charge neutrality and stabilize the structure. As expected from the structure, the material is highly anisotropic with resistive anisotropy at room temperature of 1:100:10⁵ and possesses a single, half-filled band [9]. $(\text{TMTSF})_2\text{PF}_6$ is triclinic so that the lattice vectors are not orthogonal but roughly speaking the a axis lies along the stacks (the most conducting direction), the b axis (the next most conducting direction) in the TMTSF planes and the c axis (the least conduction direction) out of the planes. At ambient pressure the material is a spin density wave insulator, but at pressures above about 6 kilobar, the ground state is superconducting. It is at such pressures and in finite magnetic fields [10] that we believe the incoherent interplane transport is realized. The theoretical picture [6,11] is that in zero field the interplane hopping is just barely sufficient to stabilize a three dimensional Fermi liquid (were it not for the superconducting transition). For a magnetic field applied along c , which minimally disrupts interplane motion, the superconductivity can be removed while retaining interplane coherence, and the behavior should be roughly Fermi liquid like. However, for fields of sufficient strength in other directions [11], the magnetic field interferes

with the interchain coherence [12] and even moderate fields in the b direction completely remove interplane coherence.

What then are the testable predictions of this theoretical description? The most natural experimental probes for coherence effects are low temperature magnetotransport measurements, and it was a number of anomalies in these measurements that first attracted our attention to the material. Consider the data depicted in Figure 1. So anisotropic a material should have a Fermi surface consisting of a pair of well separated sheets and therefore the magnetoresistance in the most conducting direction is expected in FLT to be very small and to saturate quickly [13]. Instead the material displays an enormous, angle dependent magnetoresistance for fields rotated in the bc plane and current in the a direction. Particularly striking are the dip features which occur when the magnetic field parallels a real space lattice direction. In our proposal, the dips are naturally explained as places where the magnetic field is ineffective in disrupting interplane hopping and coherence is not fully destroyed. Such a picture has a number of qualitative predictions such as the narrowing of the c direction dip roughly linearly with magnetic field and the hierarchy of appearance of the dips (the c dip must appear first as a function of field strength) [6], however, it is away from the dips - where the quantum-classical metal should be realized - that the strongest experimental consequences of the confinement of coherent motion to the planes should be apparent. In that state it is impossible to observe interference effects between interplane histories and therefore the orbital magnetoresistance contribution from the components of the magnetic field lying in the ab plane *must vanish identically*. Therefore, the magnetoresistance data away from the magic angles are also plotted in Fig. 1, not only versus angle [14], but also versus the component of the field perpendicular to the ab plane.

Note the extent to which the data away from the magic angles collapse onto a single curve, signaling the predicted independence of the magnetoresistance from in-plane field strength. The data from within the magic angle dips, where some coherence is present in our picture, are not at all independent of in-plane field strength. There is at present no other theoretical explanation for this behavior (*e.g.* all the scenarios in [15] fail to exhibit

the scaling observed in Fig. 1); at a minimum, it implies that away from the dips, and only away from the dips, the effect on the resistivity in this direction of interference effects due to fields in the ab plane is nearly zero. Since there might be reasons for this other than the incoherence of interplane hopping, let us consider further possible probes of the coherence of interplane hopping.

Clearly, the most natural quantity to consider in probing c axis coherence is the conductivity in this direction. In a simple FLT model with an isotropic scattering rate the inverse of the conductivity is easily calculated and behaves as $R_0 * (1 + \tau^2 (ev_F H_b)^2)$ where H_b is the component of the magnetic field lying in the ab plane and perpendicular to a . In low fields, we expect coherence and our theory of the magnetoresistance predicts that something like the above behavior should be observed in the c direction resistivity, while in high fields one expects the magic angle dips and the resistivity away from the dips to be independent of the in-plane field strength. Recent data of Chashechkina [16] are shown in Figure 2. The crossover from the low field, approximately FLT behavior to magic angle behavior is striking. Again, away from the magic angle dips, the resistivity depends only on the out of plane component of the field, which is now the field nearly parallel to the current! At a minimum, the data require that away from the dips, and only away from the dips, the effect of interference effects, due to fields *in* the ab plane, on the resistivity *out* of the ab plane, is nearly zero. Again, there is no theoretical proposal other than the incoherence of the interplane hopping which can account for the observed effect, especially given the qualitative agreement between the low field data and the expected behavior of a relatively clean Fermi liquid.

Danner [17] has also examined another interference effect [18] which is sensitive to the coherence of c axis transport and the existence of three dimensional Fermi surface in $(\text{TMTSF})_2\text{PF}_6$. The results are again naturally explained by the presence of coherence in fields whose projection along certain real space lattice vectors is sufficiently small, but the *total absence* of interplane coherence for other fields.

It is, of course, impossible to demonstrate the total absence of coherence experimentally:

one can only show that various measured quantities are consistent with the absence of interplane coherence, i.e. they exhibit no signs of interference effects between histories involving interplane motion. Consequently, the above results cannot be taken to *prove* the incoherence of c -axis electronic motion; however, it is truly remarkable that three different measures of interplane coherence show the complete absence of such interference effects. The magnetoresistance out of the ab plane is particularly striking since it most directly probes the c axis charge transport and, while its low field behavior fits well into the expectations of FLT, its behavior in high fields is totally anomalous from the FLT point of view.

If one accepts that there is no coherent interplane motion, is there any other explanation besides our proposal of relevant interplane hopping which has been driven incoherent by in-plane interaction effects? Clearly, if the hopping were irrelevant, *i.e.* the effective low energy theory describing the system had no out of plane hopping, the unimportance of magnetic fields in the ab plane would follow naturally. What would not make sense, however, would be the strong angular dependence of the c direction magnetoresistance (particularly the dips) or even the mere existence of a substantial, low temperature c axis conductivity. Further, if the low energy theory has no c axis hopping of charge carriers, then the resistivity must diverge as $T \rightarrow 0$. The experimental behavior in the incoherent phase is shown in Fig. 3 and there is no evidence for such behavior down to below a Kelvin.

Moreover, measurements at 50 mK (Fig. 3) show no signs of a diverging resistivity. It appears then that an explanation of the unimportance of magnetic fields in the ab plane based on the absence of c axis hopping is untenable. If the hopping exists and is incoherent, is this really a property of a low energy fixed point or is it the result of inelastic scattering or of disorder? The essential features of the magnetoresistance remain down to at least 50 mK (Fig. 3) so that an explanation based on inelastic effects again appears untenable. For many reasons, inelastic, disorder scattering is also unable to account for the loss of coherence [19] since: (1) the magic angles and scaling are only observed in the highest quality crystals; (2) the low field data are consistent with a FLT type behavior with a scattering rate of about $1K$, a rate too low to credibly explain the incoherence of c axis transport; (3) the magic

angle dips in themselves demonstrate a strong dependence of the in-plane transport on the coherence of the c axis motion, a dependence which is impossible if disorder dominates c axis motion.

VI. CONCLUSION

It therefore appears to us that the experimental situation is remarkably compelling. All indications are that in the limit of a pure system and zero temperature, $(\text{TMTSF})_2\text{PF}_6$ exhibits a phase in which an applied magnetic field [10] not only destroys superconductivity, but drives the system to a state of matter characterized by finite, relevant interplane electron hopping but the complete absence of observable interference effects between histories involving interplane motion. More succinctly, at zero temperature there is a non-vanishing interplanar electron conductivity, but it is completely incoherent. This is to be contrasted with the coherent in-plane transport. The state is, therefore, a non-Fermi liquid metal of a novel type, characterized by “quantum” in-plane, and “classical” inter-plane, transport: in short, a quantum-classical metal.

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REFERENCES

- [1] see, *e.g.*, W. Zurek, *Physics Today*, October 1991, and references therein; W. H. Zurek, Phys. Rev. D **26**, 1862 (1982), and references therein.
- [2] For a discussion of recent Schrödinger cat-like experiments, see W. Zurek, *Physics World* **10**, 24 (1997).
- [3] L. D. Landau, Sov. Phys. JETP **3**, 920 (1956); *ibid* **8**, 70 (1959).
- [4] F. D. M. Haldane, Phys. Rev. Lett. **47**, 1840 (1981) and references therein.
- [5] D. G. Clarke, S. P. Strong and P. W. Anderson, Phys. Rev. Lett. **72**, 3218 (1994); D. G. Clarke and S. P. Strong, J. Phys. Cond. Matt **8**, 10089 (1996); *ibid* **9**, 3853 (1997).
- [6] D. G. Clarke and S. P. Strong, Adv. in Phys., in press.
- [7] see, *e.g.*, A. J. Leggett *et al.*, Rev. Mod. Phys. **59**, 1 (1987) and references therein.
- [8] Here L is the system size and Z is the overlap between the bare electron and the nearly free quasiparticle of FLT.
- [9] The conduction band would be quarter filled but at the temperatures considered here the material is weakly dimerized along the most conducting direction, resulting in a half-filled band.
- [10] The field required for incoherence appears to be pressure dependent, decreasing with increasing pressure. It is possible that the quantum-classical metal can occur in this material in zero magnetic field under sufficient pressure.
- [11] S. P. Strong, D. G. Clarke and P. W. Anderson, Phys. Rev. Lett. **73**, 1007 (1994).
- [12] Coherence may similarly survive if the field is along a real space lattice direction with a component out of the ab plane and there is effective hopping in this direction of sufficient strength. Hence the presence of additional magic angle dips.

- [13] V. M. Yakovenko and A. Zhelezhyak, *Proceedings of the International Conference on the Science and Technology of Synthetic Metals*, Syn. Metals **69-71** (1994).
- [14] In the phase with incoherent interplane hopping the orbital component of the magnetoresistance cannot depend on the in plane components of the field; however, there should be small Zeeman contributions from these components even in the incoherent phase. Empirically, this contribution is unobservable.
- [15] P. M. Chaikin, Phys. Rev. Lett. **69**, 2831 (1992); K. Maki, Phys. Rev. B **45**, 5111 (1992); T. Osada, S. Kagoshima and N. Miura, Phys. Rev. B **46**, 1812 (1992); A. G. Lebed, J. Phys. I France **4**, 351 (1994).
- [16] E. I. Chashechkina and P. M. Chaikin, *Proc. of LT21*, Czech. J. Phys **46** Suppl. S5, 2649 (1996).
- [17] G. M. Danner and P. M. Chaikin, Phys. Rev. Lett. **75**, 4690 (1995).
- [18] G. M. Danner and P. M. Chaikin, Phys. Rev. Lett. **72**, 3714 (1994).
- [19] Disorder in the hopping between planes is an even less tenable proposal than in-plane disorder; as long as there is any spatially uniform component to $t_{\perp}(x)$, its Fourier transform has a delta function at zero momentum transfer with finite weight. This finite amplitude for momentum conserving single electron, interplane motion is generically far more relevant than any momentum non-conserving hopping process which might be invoked to account for the loss of interplane coherence and should always dominate over such processes in the low energy limit.

FIGURES

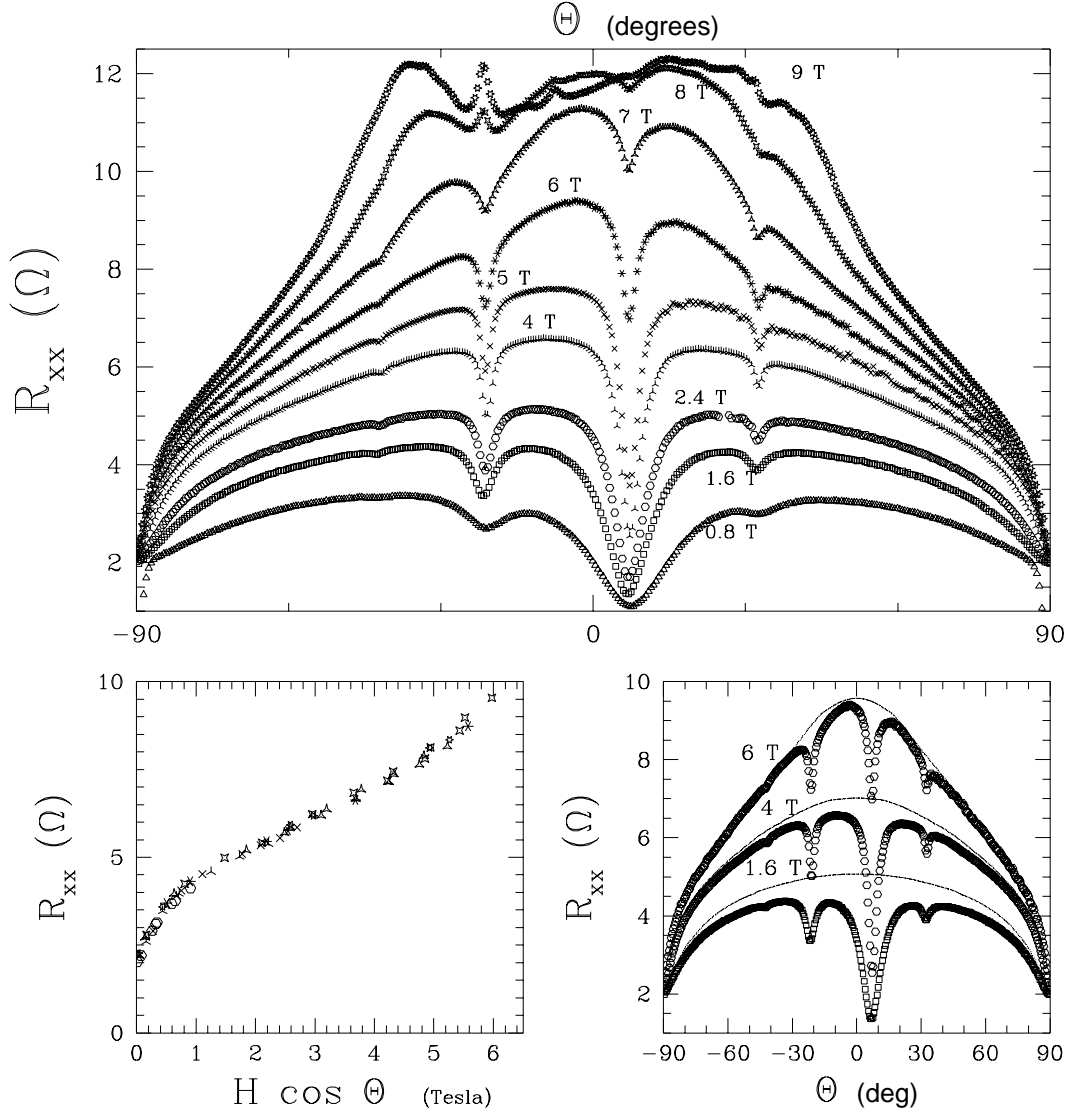


FIG. 1. **(A, Top-Center):** Magnetoresistance in the most conducting lattice direction, a , as fields of various strengths are rotated in the plane of the other two lattice directions, i.e. from $-b$ thru c on to b . Θ is measured from the perpendicular to the b axis. Data were taken at 10 kilobar and 0.5 Kelvin. **(B, Bottom-Left):** Subset of data from A for field orientations away from the “magic angles”. Data are replotted as resistance versus field strength out of the ab plane. Note the collapse onto a single scaling curve irrespective of field strength (see text). **(C, Bottom-Right):** Data for field strengths of 1.6, 4 and 6 Tesla together with the expected magnetoresistance from the scaling curve of part B. Deviations from scaling occur only within a vicinity of the magic angles that decreases rapidly with field.

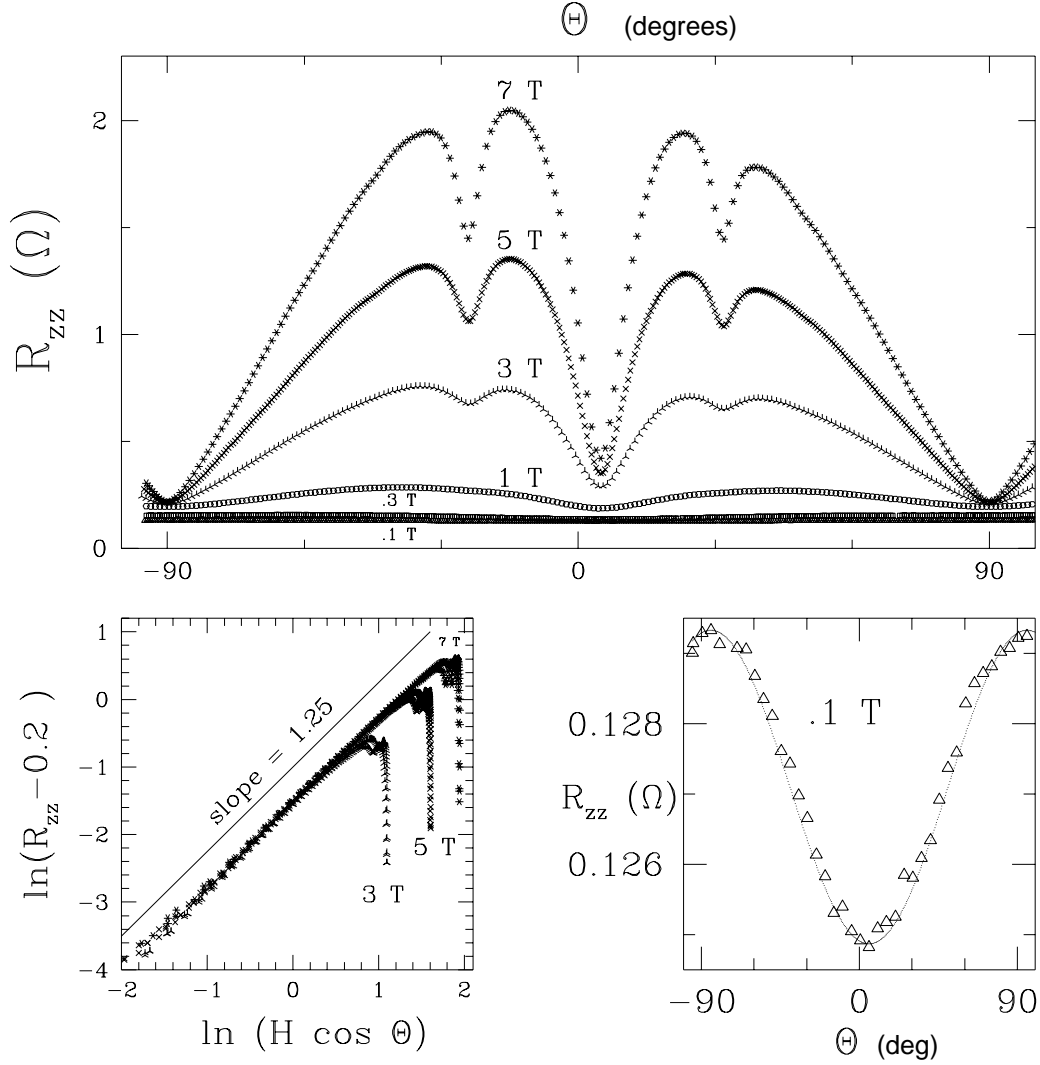


FIG. 2. **(A, Top-Center)**: Magnetoresistance perpendicular to the ab plane as fields of various strengths are rotated as in Figure 1. Data were taken at 10 kilobar and 1.3 Kelvin. **(B, Bottom-Left)**: Data from A for 3, 5, and 7 Tesla plotted as natural logarithm of deviation from a reference value versus natural logarithm of magnetic field strength perpendicular to the ab plane. Note that, away from the “magic angles”, the data exhibit power law dependence on only one component of the magnetic field: $\Delta R \propto (H \cos \Theta)^p$, $p \sim 1.25$ (see text). **(C, Bottom-Right)**: Weak field (0.1 Tesla) magnetoresistance for bc rotation. Dotted line is a fit to the data of the form $R_0 + \alpha|H \times \hat{c}|^2$, the semi-classical prediction.

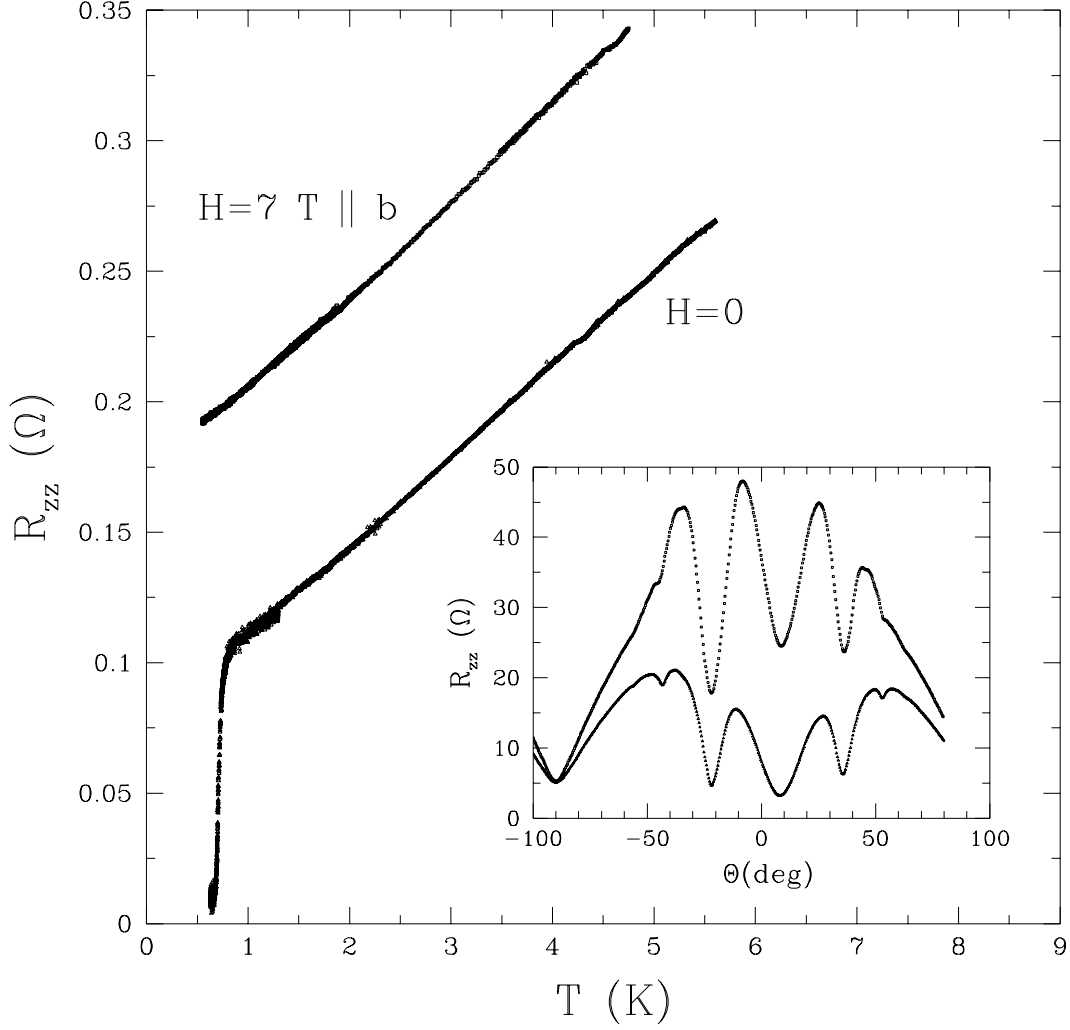


FIG. 3. Temperature dependence of the resistivity perpendicular to the ab plane both in the absence of a magnetic field and in a field of 7 Tesla along b (the latter condition is in the incoherent phase). Data taken at 10 kilobar applied pressure (see text). **Inset:** Magnetoresistance perpendicular to the ab plane for bc field rotations as in Figure 2 but at 8.2 kilobar and 50 millikelvin.